Why more hadronlike photons produce less particles on nuclear targets

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Abstract

In hadron-nucleus interactions, the stronger is nuclear shadowing in the total cross section the higher is the multiplicity of secondary hadrons. In deep inelastic scattering, nuclear shadowing at small x is associated with the hadronlike behaviour of photons as contrasted to the pointlike behaviour in the non-shadowing region of large x. In this paper we predict smaller mean multiplicity of secondary hadrons, and weaker fragmentation of the target nucleus, in deep inelastic leptoptoproduction on nuclei in the shadowing region of small x as compared to the non-shadowing region of large x. This paradoxial conclusion has its origin in nuclear enhancement of the coherent diffraction dissociation of photons. We present numerical predictions for multiproduction in μXe interactions studied by the Fermilab E665 collaboration.

1

1 Introduction

At high energies ν and/or very large $\frac{1}{x} \gg 1$, where x is the Bjorken variable, the real, and virtual, photoabsorption can conveniently be considered as interaction with the target nucleon (nucleus) of hadronic (multiparton) Fock states X of the photon (for the review and references to the early work see [1]). The best known consequence of this hadronic Fock-state mediated photoabsorption is that nuclear shadowing in the forward Compton scattering off nuceli, i.e., in the nuclear structure function, will be similar to that in the hadronic XA scattering [2]. In the framework of multiple scattering theory [3], nuclear shadowing has its origin in diffractive excitation $\gamma^* \to X$ of the intermediate state X followed by its deexcitation $X \to \gamma^*$ inside the nucleus [4]. The closely related process is the direct diffractive excitation of hadronic Fock states of the photon $\gamma^* + N(A) \to N(A) + X$, which is inseparable from nuclear shadowing. Although at large Q^2 deep inelastic scattering (DIS) will eventually be dominated by interactions of the small-size multiparton Fock states, the remarkable feature of QCD is that hadronic-size Fock components of the photon contribute significantly even at asymptotically large Q^2, and completely dominate diffraction dissociation of photons [5,6]. In this paper we discuss a novel, and paradoxial, feature of diffractive DIS on nuclei: weaker multiproduction of secondary hadrons, and weaker fragmentation of the target nucleus, in the shadowing region of $x \ll 1$, when photons become hadronlike, with respect to the multiproduction in the nonshadowing region of large x, where photons are pointlike.

Because of nuclear shadowing, the nuclear cross section is smaller than A times the free-nucleon cross section (hereafter A is the nuclear mass number). In the multiple scattering theory [3], this shadowing in the hadron-nucleus interaction comes from multiple intranuclear rescatterings of the projectile. A convenient parameter which measures the strength of multiple rescatterings is (for the review see [1])

$$\bar{\nu} = \frac{A\sigma_N}{\sigma_A}.\tag{1}$$

Multiple intranuclear interactions are one of the sources of stronger multiproduction on nuclei, and a good empirical approximation for the scaled nuclear multiplicity of secondary particles $R = \langle n_A \rangle / \langle n_N \rangle$ is

$$R \approx \frac{1}{2}(1 + \bar{\nu})\,. \tag{2}$$

The larger is the free-nucleon cross section σ_{hN} the stronger is the nuclear shadowing (for the review see [1,7]). For instance, because $\sigma_{pN} > \sigma_{\pi N} > \sigma_{KN}$, one has $\bar{\nu}_{pA} > \bar{\nu}_{\pi A} > \bar{\nu}_{KA}$, and indeed experimentally the scaled nuclear multiplicity satisfies $R_{pA} > R_{\pi A} > R_{KA}$ (for the review see [1]).

For the weakly interacting pointlike probes, like in DIS in the non-shadowing (NS) region, the impulse approximation is exact and the amplitude of forward elastic scattering on a nucleus equals A times the free-nucleon amplitude:

$$F_A = A F_N\,. \tag{3}$$

This implies $\bar{\nu} = 1$. Furthermore, because by virtue of the optical theorem, the n-particle production cross section is related to the n-particle discontinutity of the forward scattering amplitude. Then, if taken at face value, Eq. (3) would have imlpied identical discontinuities of the nuclear and the free-nucleon amplitudes, i.e., identical multiplicities for the nuclear and free-nucleon interactions. This naive expectation fails, though, because of the cascading effects [1,2,8-12], which do not affect the nuclear cross section, but contribute to the particle production. The significance of cascading as a necessary condition for thermalization of the produced particles and for the formation of the quark-gluon plasma in collisions of ultrarelativistic heavy ions is discussed in [12]. There is a mounting experimental evidence for cascading in leptoproduction on nuclei [12,13].

It was suggested quite a time ago [1,2,8,9], that in the leptoproduction on nuclei, one can control $\bar{\nu}$ by varying the Bjorken variable x from the non-shadowing (NS) region of large $x \gtrsim 0.05$, where $\bar{\nu} = 1$, to the shadowing (SH) region of very small x, where $\bar{\nu} > 1$. For instance, in the μXe scattering at $x \sim 10^{-3}$ the shadowing effect is rather strong, $\bar{\nu} \sim 1.5$ [14-16]. Then, the empirical law Eq. (2) would have suggested a strong, $\sim 25\%$, nuclear enhancement of the mean multiplicity. Apart from the larger mean multiplicity of secondary paricles, in the shadowing region of small x one would naively have expected other signals

of enhanced intranuclear reinteractions like the higher multiplicity of knocked-out protons (grey tracks) and of the nucleus fragmentation in general, which rise with $\bar{\nu}$ [1,2,8,9].

In this communication we wish to demonstrate how the above summarized conventional wisdom fails: the very mechanism of the hadronlike behavior of (virtual) photons leads to a weaker multiproduction on, and weaker fragmentation of, nuclei in the shadowing region as compared to multiproduction in the non-shadowing region. The principal observation goes as follows: In the free-nucleon interactions, the diffraction dissociation (DD) events a characterized by a large (pseudo)rapidity gap (LRG) between the recoil proton and the hadronic debris from the diffraction dissociation of photons. Because of this large rapidity gap, the DD events have smaller mean multiplicity than the non-diffraction dissociation (ND) events. The major finding of the present paper is that, in DIS on nuclei, the fraction of DD and/or LRG events significantly rises with A. On the black-disc nucleus, the coherent DD, which leaves the target nucleus in the ground state and consequently gives a vanishing hadronic activity in the nucleus fragmentation region, will make $\sim 50\%$ of the total DIS cross section. Nuclear shadowing and DD come in one package, and because of nuclear enhancement of DD the hadronlike photons produce less secondary hadrons than the pointlike photons. Recently, there was much theoretical interest in DD of photons [5,6], and LRG events were observed in DIS at HERA [17] with the rate which agrees with the theoretical prediction [5]. The novel manifestation of DD of photons, discussed in this paper, adds to the growing interest in the large-rapidity gap physics in DIS.

The observation of different A-dependence of the diffractive and nondiffractive multiproduction in the hadron-nuclei collisions was made by one of the authors quite a time ago [18]. The major difference between the leptoproduction and hadroproduction is that in the latter case the nuclear DD only makes a small fraction of the nuclear cross section, whereas in the leptoproduction DD cross section is much larger.

The paper is organized as follows. In section 2 we start with the brief review of the dipole-cross section approach to diffractive DIS. In section 3 we derive the cross section for the coherent and incoherent DD on nuclei and discuss the relationship between the DD of photons and nuclear shadowing. We also demonstrate the nuclear enhancement of DD

cross section. In section 4 we discuss the impact of finite energy effects on the so-called triple-pomeron component of the coherent and incoherent DD on nuclei and on nuclear shadowing, and present predictions for the A dependence of the rapidity gap distribution. The experimental signatures of nuclear enhancement of diffraction dissociation are discussed in section 5. In section 6 we comment on why the effects of DD in the leptoproduction and hadroproduction on nuclei are so much different. Our principal results and conlcusions are summarized in section 7.

2   Hadronic prioperties of the photon and the dipole-cross section representation

We start with the brief review of the dipole-cross section representation for diffractive DDIS [14,5,6], which provides a unified description of nuclear shadowing and of diffraction dissociation of photons. At small x, DIS can be viewed as interaction of the hadronic fluctuations the virtual photon transforms into at large distance

$$\Delta z \sim \frac{2\nu}{Q^2 + M^2} \sim \frac{1}{m_p x} \gtrsim R_N, R_A \qquad (4)$$

in front of the target nucleon (nucleus) [1]. Here ν and Q^2 are the laboratrory energy and virtuality of the photon, M is the invariant mass of hadronic fluctuation of the photon and $R_{A,N}$ is the radius of the target nucleus (nucleon). Because of $\Delta z \gtrsim R_A$, the transverse separation $\vec{r}$ of partons in the multiparton Fock state of the photon becomes as good a conserved quantity as the angular momentum. The resulting diagonalization of the diffractive S-matrix in the $\vec{r}$-representation leads to a very simple, and intuitively appealing, description of diffractive interactions in the dipole-cross section approach.

We present the approach starting with interactions of the simplest $q\bar{q}$ Fock state of the photon. The principal quantities are the total cross section $\sigma(r)$ for interaction of the colour dipole, i.e., the colour-singlet $q\bar{q}$ pair with the transvrese separation $\vec{r}$ with the nucleon target, and the wave functions $|\Psi_{T,L}(\alpha, \vec{r})|^2$ for the (T) transverse and (L) longitudinal photons, computed in [14]. Here α is a fraction of the lightcone momentum of the photon

carried by the quark of the $q\bar{q}$ pair. The total photoabsorption cross section and the inclusive forward DD cross section for the free-nucleon target are given by

$$\sigma_{T,L} = \int_0^1 d\alpha \int d^2\vec{r} \, |\Psi_{T,L}(\alpha, \vec{r})|^2 \sigma(r) = \langle \sigma(r) \rangle_{T,L} \tag{5}$$

$$\left. \frac{d\sigma_D^{(N)}}{dt} \right|_{t=0} = \int dM^2 \left. \frac{d\sigma_D^{(N)}}{dM^2 dt} \right|_{t=0} = \int_0^1 d\alpha \int d^2\vec{r} \, |\Psi_{T,L}(\alpha, \vec{r})|^2 \frac{\sigma(r)^2}{16\pi} = \frac{1}{16\pi} \left\langle \sigma(r)^2 \right\rangle_{T,L} \tag{6}$$

In the diffraction production of the state of mass M the target proton receives a small recoil momentum κ which in the laboratory frame equals

$$\kappa = \frac{M^2 + Q^2}{2m_N\nu} = m_N x(1 + \frac{M^2}{Q^2}) = m_N x[1 + \exp(y)] \,, \tag{7}$$

Here

$$y = \log\left(\frac{M^2}{Q^2}\right) \tag{8}$$

is a convenient variable which measures the mass of the diffractively excited state. As we shall see below, DD has simple scaling properties in terms of this variable y, and we strongly advocate an analysis of DD in terms of this new variable. Let W be the total collision energy in the photon-proton c.m.s, $W^2 = 2m_p\nu - Q^2$. The recoil proton emerges in the final state separated from the hadronic debris of the photon by large (pseudo)rapidity gap

$$\Delta\eta \approx \log\left(\frac{W^2}{M^2}\right) = \log(\frac{1}{x}) - \log(\frac{M^2}{Q^2}) = \log(\frac{1}{x}) - y \tag{9}$$

For the reaction to be the diffraction dissociation, the (pseudo)rapidity gap $\Delta\eta$ must be large, $\Delta\eta \gtrsim 2.5 - 3$ (for the recent review on diffraction dissociation in hadronic scattering see [19]). The total (pseudo)rapidity span equals

$$Y_{max} \approx \log(\frac{1}{x}) + \log(\frac{Q^2}{\langle p_\perp \rangle^2}) \,, \tag{10}$$

where $\langle p_\perp \rangle$ is the mean transvrese momentum of secondary hadrons, $\langle p_\perp \rangle \sim \frac{1}{2}m_\rho$. The maximal kinematically allowed rapidity gap $\Delta\eta \sim Y_{max}$ corresponds to exclusive production of the very low-mass state like the continuum two-pion state near the threshold and/or the ρ^0 meson, $M_{min} \sim \frac{1}{2}m_\rho$, and in DIS

$$y_{min} = \log\left(\frac{M_{min}^2}{Q^2}\right) < 0 \,. \tag{11}$$

6

Notice, that $y = 0$ corresponds to the rapidity of the virtual photon.

Hereafter we concentrate on the dominant DD of transverese photons and suppress the subscript T. Excitation of $q\bar{q}$ pairs leads to the mass spectrum peaked at $M^2 \sim Q^2$ [5]:

$$\left.\frac{d\sigma_D^{(N)}}{dtdM^2}\right|_{t=0} \approx \Sigma_{DD}\frac{M^2}{(Q^2 + M^2)^3}\,. \qquad (12)$$

DD of photons can also be viewed as DIS on pomerons ($\mathbb{IP}$), and the mass spectrum can be related to the pomeron structure function. Diffraction excitation of the $q\bar{q}$ Fock state of the photon corresponds to DIS on the valence $q\bar{q}$ component of the pomeron, with the structure function [5,6] $F_2^{\mathbb{IP}}(\beta) \propto \beta(1 - \beta)$, where $\beta = Q^2/(Q^2 + M^2)$ is the Bjorken variable for the $e\mathbb{IP}$ deep inelastic scattering. This is the smooth spectrum, which does not contain an explicit ρ^0 resonance contribution, but it correctly reproduces the resonance-smeared mass spectrum even in the limit of real photoproduction $Q^2 = 0$ [5,20]. The normalization of the mass spectrum Σ_{DD} is such that the integrated diffraction excitation cross section

$$\sigma_D^{(N)} = r_D\sigma_N = \int dM^2 dt \left.\frac{d\sigma_D}{dtdM^2}\right|_{t=0} \approx \int dM^2 \frac{1}{b_D}\left.\frac{d\sigma_D}{dtdM^2}\right|_{t=0} \qquad (13)$$

makes the fraction $r_D \sim 10\%$ of the total photoabsorption on free nucleons. This fraction r_D was predicted to only weakly depend on x and Q^2 [5,21], which is in good agreement with the first experimental data on DD of photons from the ZEUS collaboration at HERA [17]. The above estimate $r_D \sim 10\%$ for excitation of $q\bar{q}$ Fock states of the photon is found [5,21] with the dipole cross section $\sigma(r)$ of Ref. [14] and for the diffraction slope $b_D \sim b_{\pi N} \sim 10\text{GeV}^{-2}$. The effects of higher $q\bar{q}g, ...$ Fock states will be discussed below.

Let $n(W^2)$ be the mean multiplicity in the generic inelastic interaction. Then, in the DD events the mean multiplicity will be approximately equal to

$$n_D(M^2) \approx n(W^2 \exp(-\Delta\eta)) < n(W^2)\,. \qquad (14)$$

Consequently, the enhancement of the fraction of DD events in DIS on nuclei results in the lower mean multiplicity of secondary particles in the diffractive multiproduction on nuclei. In the next section we shall demonstrate that such an enhancement indeed takes place. Another important signature of DD is a very small recoil of the target nucleon (nucleus), which means a lack of any observable hadronic activity in the target region.

3 Hadronic properties of the photon and deep inelastic scattering off nuclei

In this section we discuss the impact of diffraction dissociation in the leptoproduction off nuclei, starting with the high-energy limit of $\Delta z \gg R_A$. In the interaction with nuclear targets one has to distinguish the (coh) coherent diffraction dissociation $\gamma^* + A \to X + A$, when the target nucleus remains in the ground state, and the (inc) incoherent diffraction dissociation $\gamma^* + A \to X + A^*$, when one summs over all excitations and breakup of the traget nucleus not followed by the secondary particle production in the nucleus fragmentation region. The both processes lead to the LRG events, but have different A-dependence and slightly different dependence on the rapidity gap $\Delta \eta$.

We derive the cross sections for the coherent and incoherent diffraction dissociation on a nucleus using the technique developed in [22,23]. We start with the total cross section of photoabsorption on a nucleus which equals

$$\sigma_A = R_{sh} A \sigma_N = A \sigma_N - \Delta \sigma_{sh} = 2 \left\langle \int d^2\vec{b} \{ 1 - \exp[-\frac{1}{2}\sigma(r)T(b)] \} \right\rangle$$

$$= A \langle \sigma(r) \rangle - \int d^2\vec{b} \left\langle \exp[-\frac{1}{2}\sigma(r)T(b)]\{1 - [1 - \frac{1}{2}\sigma(r)T(b)] \exp[\frac{1}{2}\sigma(r)T(b)]\} \right\rangle$$

$$= A \langle \sigma(r) \rangle - \frac{1}{4} \int d^2\vec{b} T(b)^2 \left\langle \sigma(r)^2 \exp[-\frac{1}{2}\sigma(r)T(b)] \right\rangle + ... \qquad (15)$$

Here $\vec{b}$ is the impact parameter, $T(b) = \int dz n_A(z, \vec{b})$ is the optical thickness of the nucleus and $n_A(z, \vec{b})$ is the nuclear matter density (for the nuclear density parametrizations see [24]). In Eq. (15) we decomposed the nuclear cross section into the impulse approximation term $A\sigma_N$ and the shadowing term $\Delta\sigma_{sh}$. Although the cross section of photoabsorption on the free nucleon is small,

$$\sigma_N = \frac{4\pi^2 \alpha_{em}}{Q^2} F_2(x, Q^2), \qquad (16)$$

and rapidly vanishes with rising Q^2, the nuclear shadowing persists at all Q^2 and the ratio $\Delta\sigma_{sh}/\sigma_A = 1 - R_{sh}$ is approximately independent of Q^2 [14]. The leading term of the shadowing is shown in the last line of Eq. (15). For the light nuclei, it rises with the atomic number $\propto A^{1/3}$:

$$1 - R_{sh} \propto \frac{1}{A} \int d^2\vec{b} T(b)^2 \propto \frac{A}{R_A^2} \propto A^{1/3}. \qquad (17)$$

For the heavier nuclei, this rise slows down because of the nuclear attenuation factor in the integrand of (15). Evidently, the driving term of the shadowing cross section is proportional to the forward DD on free nucleons Eq. (6) [4,14,25], and the experimental observation of nuclear shadowing already is a solid evidence for significant DD of virtual photons in DIS.

The nuclear profile function for the coherent DD $\gamma^* + A \to X + A$ equals

$$\Gamma_{coh}(\gamma^* \to X, \vec{b}) = \langle X | \{1 - \exp[-\frac{1}{2}\sigma(r)T(b)]\} | \gamma^* \rangle. \tag{18}$$

The total cross section of the coherent DD on a nucleus, integrated over all final states X, equals

$$\sigma_{coh} = R_{coh} A \sigma_D^{(N)} = \int d^2\vec{b} \sum_X |\Gamma_{coh}(\gamma^* \to X, \vec{b})|^2$$

$$= \int d^2\vec{b} \left\langle \{1 - \exp[-\frac{1}{2}\sigma(r)T(b)]\}^2 \right\rangle = \frac{1}{4} \int d^2\vec{b} T(b)^2 \left\langle \sigma(r)^2 \exp[-\sigma(r)T(b)] \right\rangle + ... \tag{19}$$

Here the closure was used, and in the last line of Eq. (19) we show the leading term of the coherent DD cross section, which has a close semblance to the leading term of the nuclear shadowing in Eq. (15).

The differential cross section of the incoherent production $\gamma^* + A \to X + A$ equals

$$\frac{d\sigma_{inc}(\gamma^* \to X)}{dt}\bigg|_{t=0} = R_{inc} A \frac{d\sigma_D^{(N)}}{dt}\bigg|_{t=0}$$

$$= \frac{1}{16\pi} \int d^2\vec{b} T(b) \left| \langle X | \sigma(r) \exp[-\frac{1}{2}\sigma(r)T(b)] | \gamma^* \rangle \right|^2 + ... \tag{20}$$

Here we only have shown the single incoherent scattering term, which is sufficient for the practical purposes, the full multiple-scattering expansion can be found in [22]. Summing over all produced states X making use of the closure, we can write

$$\sigma_{inc} = R_{inc} A \sigma_D^{(N)} = \sigma_D^{(N)} \int d^2\vec{b} T(b) \frac{\langle \sigma(r)^2 \exp[-\sigma(r)T(b)] \rangle}{\langle \sigma(r)^2 \rangle} \tag{21}$$

Diffraction excitation of $q\bar{q}$ pairs describes only a part of nuclear shadowing and of the DD cross section, but this part is the dominant one at moderate energies and/or large rapidity gaps. Furthermore, the A-dependence of diffraction excitation of higher Fock states of the photon will be similar to that for the $q\bar{q}$ Fock states. In Fig. 1 we present our predictions for the A-dependence of the normalized coherent R_{coh} and the incoherent R_{inc} diffraction

dissociation of photons into $q\bar{q}$ pairs. We also show the contribution of excitation of $q\bar{q}$ pairs to nuclear shadowing. The calculations are based on the dipole cross section of Ref. [14], which has already been successfully applied to the quantitative description of nuclear shadowing in DIS [14,15] and of colour transparency effects in exclusive leptoproduction of the ρ^0 mesons on nuclei [25].

The incoherent and coherent DD have quite a different A-dependence. In the limit of vanishing nuclear attenuation, when $\exp[-\sigma(r)T(b)] \approx 1$, we have $\int d^2\vec{b}\,T(b)\exp[-\sigma(r)T(b)] = A$, and for light nuclei $\sigma_{inc} \approx A\sigma_D$ and $R_{inc} \approx 1$. For heavier nuclei, when the attenuation factor $\exp[-\sigma(r)T(b)]$ in the integrand of (21) becomes important, the relative fraction of the incoherent production decreases with A, i.e., $R_{inc} < 1$. As a matter of fact, attenuation is substantial already for the carbon nucleus, and for very heavy nuclei $R_{inc} \ll 1$. In the opposite to that, in the limit of weak attenuation for light nuclei, the relative fraction of the coherent diffraction dissociation rises $R_{coh} \propto A^{1/3}$, see Eq. (17), and the coherent production very rapidly takes over the incoherent production. In the lightest nuclei (deuterium,....) the coherent cross section is smaller than the incoherent one. For the heavier nuclei this rise of the coherent cross section slows down because of the attenuation factor $\exp[-\sigma(r)T(b)]$ in the integrand of (19), but persists for the whole range of nuclei. Furthermore, for very heavy nuclei, in the black-disc limit, we predict that the coherent DD makes one half of the total DIS cross section, see the discussion in Section 6. A comparison of Eqs. (19),(15) shows that for the light nuclei the coherent production cross section and the nuclear shadowing cross section are very close to each other:

$$R_{coh/Sh} = \frac{\sigma_{coh}}{\Delta\sigma_{sh}} = \frac{R_{coh}}{1 - R_{sh}} \approx 1\,. \tag{22}$$

For the heavier nuclei this ratio $R_{coh/Sh}$ slowly decreases because of stronger nuclear attenuation factor $\exp[-\sigma(r)T(b)]$ in the coherent production cross section σ_{coh} Eq. (19) as compared to the attenuation factor $\exp[-\frac{1}{2}\sigma(r)T(b)]$ in the shadowing cross section $\Delta\sigma_{sh}$ Eq. (15). In Fig. 1 we show the ratio $R_{D/Sh}$ of the total, coherent plus incoherent, DD cross section to the shadowing cross section:

$$R_{D/Sh} = \frac{\sigma_{coh} + \sigma_{inc}}{\Delta\sigma_{sh}} = \frac{R_{coh} + R_{inc}}{1 - R_{sh}}\,. \tag{23}$$

For light nuclei it is slightly above unity for the sizable contribution of the incoherent DD. The results shown in Fig. 1 confirm our anticipation of the enhancement of diffraction production on nuclei with respect to the free-nucleon target.

4 The energy dependence of diffraction and the triple-pomeron component of diffraction dissociation

At high energy, when $\Delta z \gg R_A$ and the frozen-size approximation holds, the diffraction excitation of $q\bar{q}$ pairs gives the energy and/or x independent DD cross section. There are two sources of the energy dependence of DD both on nucleons and nuclei:

Firstly, at finite energy, in the diffraction excitation $\gamma^* + A \to X + A$ the target nucleus receives the longitudinal momentum transfer κ Eq. (7). Henceforce, if we ask for the coherent diffraction excitation when the target nucleus remains in the ground state, then the production amplitude will be proportional to the so-called body form factor of the nucleus $G_A(\kappa^2)$, which for all the practical purposes can be taken equal to the charge form factor of the nucleus. In the high-energy limit, which in DIS is the limit of $x \to 0$, the longitudinal momentum transfer $\kappa \to 0$ and $G_A(\kappa^2) \to 1$.

At finite energy ν and/or finite x, the effect of this form factor can be included as follows. Let $d\sigma_D^{(N)}/dM^2$ be the mass spectrum of diffraction dissociation on the free nucleon, in which case the form factor effect can be neglected for the small size of the nucleon. For the purposes of the present analysis we can write

$$\frac{d\sigma_{coh}}{dM^2} \approx R_{coh} A \frac{d\sigma_D^{(N)}}{dM^2} G_A(2\kappa^2) \tag{24}$$

and

$$\sigma_{coh} = R_{coh} A \int dM^2 \frac{d\sigma_D^{(N)}}{dM^2} G_A(2\kappa^2) \,, \tag{25}$$

where the approximation $G_A^2(\kappa^2) \approx G_A(2\kappa^2)$, which is exact for the Gaussian form factor, was made. Because of this suppression by the nuclear form factor, the coherent DD only takes place at a sufficiently small x such that $\kappa R_A \lesssim 1$, i.e.,

$$x \lesssim \frac{1}{R_A m_N} \sim 0.1 A^{-1/3} \,. \tag{26}$$

The above simplifying approximation that nuclear attenuation does not depend on the mass M of the excited state is viable for the inclusive DD cross section (for discussion of possible difference of nuclear attenuation for the production of specific exclusive final states see [26]. The above form factor suppression is absent in the incoherent DD on nuclei.

Nuclear shadowing has its origin in the destructive interference of the single-scattering (impusle-approximation) amplitude with the amplitude of the double (and higher order) scattering $\gamma^* \to X \to \gamma^*$ [4,14]. The driving term of nuclear shadowing in Eq. (15) is proportional to $d\sigma_D^{(N)}/dt|_{t=0}$. The principal finite-energy modification of Eq. (15) is as follows: The intermediate state X acquires the phase $\kappa(z_2 - z_1)$ during its propagation from the production point z_1 to the reabsorption point z_2. The corresponding contribution to the double-scattering amplitude enters with the phase factor $\exp[i\kappa(z_2 - z_1)]$ with respect to the impulse approximation amplitude. After integration over $z_{1,2}$, this phase factor gives rise to the suppression factor $G_A^2(\kappa^2) \approx G_A(2\kappa^2)$ [4,25,27]. As a result, nuclear shadowing will be proportional to

$$\Delta\sigma_{sh} \propto \int dM^2 \left. \frac{d\sigma_D^{(N)}}{dtdM^2}\right|_{t=0} G_A(2\kappa^2) \,. \tag{27}$$

Consequently, the x dependence of nuclear shadowing Eq. (27) and of the total cross section of the coherent DD Eq. (25) will be approximately the same. One must bear in mind, though, that the magnitude of the observed nuclear shadowing is somewhat reduced (by $\sim 5\%$ for heavy nuclei) because of the low-x manifestation of the nuclear EMC effect [14,15].

Second source of the energy dependence of DD and of nuclear shadowing is the triple-pomeron component of diffraction dissociation. Namely, the diffraction excitation of the $q\bar{q}$ Fock states of the photon gives the mass spectrum (12) which converges rapidly at $M^2 \gtrsim Q^2$. It is the counterpart of diffraction excitation of resonances in hadronic interactions. Diffraction excitation of the $q\bar{q}g$ and higher Fock states generates the so-called triple pomeron mass spectrum [14,5,6]

$$\frac{1}{\sigma_N} \left. \frac{d\sigma_D^{(N)}}{dtdM^2}\right|_{t=0} = \left[\frac{M^2}{M^2 + Q^2}\right]^2 \cdot \frac{A_{3\mathbb{P}}}{M^2 + Q^2} \approx \frac{A_{3\mathbb{P}}}{M^2 + Q^2} \,. \tag{28}$$

In terms of the structure function of the pomeron, diffraction excitation of the $q\bar{q}g$ (and higher) Fock states of the photon corresponds to DIS on the $q\bar{q}$ sea in the pomeron, the

sea being generated from the valense gg component of the pomeron. The particular mass spectrum (28), with the factor $[M^2/(M^2 + Q^2)]^2$ which suppresses the triple-pomeron contribution at $M^2 \lesssim Q^2$, reflects the large-β behaviour of the gluon structure function in the pomeron $G^{\mathbb{P}}(\beta) \propto (1 - \beta)$ [5,6].

Although the triple-pomeron coupling $A_{3\mathbb{P}}$ is the dimensionfull constant, it must only weakly depend on Q^2 [14,6,28] and can be borrowed from the triple-pomeron phenomenology of the real photoproduction [29]: $A_{3\mathbb{P}} \approx 0.16\text{GeV}^{-2}$. This choice of $A_{3\mathbb{P}}$ leads to an excellent quantitative description of the experimental data on nuclear shadowing [14,15]. For the reference purposes, in Fig. 2 we present our estimate for the nuclear shadowing in μXe scattering. The x-dependence of $R_{sh}(x)$ at small x is dominated by the rising contribution to shadowing from the triple-pomeron component of the mass spectrum. At $x \gtrsim 0.01$ the x-dependence of shadowing comes predominatly from the form factor effects Eq. (27). (Compared to the more detailed analysis in Ref. [15], here we neglect corrections for the nuclear EMC effect, which may reduce nuclear shadowing by $\sim 5\%$. Also, here we use simple parametrizations (12), (28) rather than the direct calculation of the mass spectrum).

In the triple-pomeron regime, the diffraction slope is smaller than in the resonance and/or $q\bar{q}$ excitation region (for instance, see [19,29]). Therefore, in our estimates of the DD cross section we take $b_{3\mathbb{P}} \approx \frac{1}{2}b_{\pi N} \approx (5-6)\text{GeV}^{-2}$.

In principle, the triple-pomeron coupling is calculable in terms of the cross of interaction of the $q\bar{q}g$ Fock state [6], and such a calculation is in progress. The same three-parton cross section controls the nuclear attenuation, which for the $q\bar{q}g$ (and higher) Fock states can be slightly different from that for the $q\bar{q}$ Fock state. At very small values of x the triple-pomeron component of shadowing takes over, but for the moderately small x of the present muon experiments the dominant contribution to the nuclear shadowing and to the diffraction dissociation cross section comes from the $q\bar{q}$ states, see Fig. 2 in which we decompose nuclear shadowing into the $q\bar{q}$ excitation and triple-pomeron components. For this reason, for the purposes of the present analysis, we can make a simplifying assumption of similar nuclear attenuation of the $q\bar{q}$ and $q\bar{q}g$ states.

Now we are in the position to write down the total mass spectrum in the diffraction

dissociation of photons on a free nucleon:

$$F_N(y) = \frac{1}{\sigma_N} \frac{d\sigma_D^{(N)}}{dy} = (M^2 + Q^2) \frac{d\sigma_D^{(N)}}{\sigma_N dM^2} =$$

$$2r_D \frac{Q^2 M^2}{(Q^2 + M^2)^2} + \frac{A_{3\mathbb{P}}}{b_{3\mathbb{P}}} \frac{M^2}{Q^2 + M^2} = 2r_D \frac{\exp(y)}{[1 + \exp(y)]^2} + \frac{A_{3\mathbb{P}}}{b_{3\mathbb{P}}} \frac{\exp(y)}{1 + \exp(y)}. \tag{29}$$

Here the normalization $2r_D$ of the $q\bar{q}$ term is written assuming that Q^2 is large enough, $Q^2 \gg M_{min}^2$. For the free-nucleon target, the differential probability of diffraction excitation of heavy masses $M^2 \gg Q^2$ flattens at large y at the value

$$\frac{d\sigma_D^{(N)}}{\sigma_N dy} = \frac{A_{3\mathbb{P}}}{b_{3\mathbb{P}}} \approx 0.025 - 0.03. \tag{30}$$

The triple-pomeron component of the mass-spectrum only becomes the dominant one at [5]

$$M^2 \gtrsim (5 - 7) Q^2 \tag{31}$$

The y-distribution $F_N(y)$ for $Q^2 = 1\text{GeV}^2$ and $x = 2 \cdot 10^{-3}$, relevant to the E665 experiment [13], is shown in Fig. 3. It is slightly peaked at $y \sim 0.5$ and exhibits the onset of the triple-pomeron plateau at large $y \gtrsim 2$, see estimate (31). Excitation of the $q\bar{q}$ pairs dominates at smaller y.

We advocate studying the rapidity gap distribution in terms of the variable y, because for the free-nucleon target, and also for the nuclear target at $x \ll 1$, the y distribution must be a scaling function of y, which depends neither on Q^2 nor x. For the comparison, in Fig.4 we show the predicted M^2-distribution for few values of x assuming $W^2 = 400\text{GeV}^2$ which is appropriate for the E665 experiment. In the mass spectrum, this nice y-scaling property is completely obscured.

In the comparison with experimental mass spectra taken at different values of x, one only must bear in mind, that Eq. (29) only holds at the (pseudo)rapidity gap $\Delta\eta \gtrsim 2.5 - 3$, because at smaller rapidity gaps the non-diffractive mechanisms of the rapidity gap generation will take over [30]. This puts a restriction on the excited mass $M^2 \lesssim 20\text{GeV}^2$ for the $W^2 \approx 400\text{GeV}^2$ of the E665 experiment. At asymptotically large W^2, all the curves must flatten at the same asymptotic value Eq. (30) at large M^2. As Eq. (9) shows with increasing Q^2 and increasing x at the fixed value of W^2, the less and less room will be left for the large-y triple-pomeron component and, at a sufficiently large Q^2, DD will be dominated by

14

excitation of the $q\bar{q}$ state. This can also be seen from Fig. 2, in which we show separately the contribution to nuclear shadowing from the triple-pomeron component.

5 Experimental signatures of nuclear enhancement of diffraction dissociation

5.1 The mass spectrum in DD on nuclei and the rapidity-gap distribution

The nuclear enhancement/attenuation and the nuclear form factor effects lead to significant changes in the rapidity gap distribution:

$$F_A(y) = (M^2 + Q^2)\frac{d\sigma_D^{(A)}}{\sigma_A dM^2} =$$

$$\frac{R_{inc} + R_{coh}G_A(2\kappa^2)}{R_{sh}(x)} \cdot \left\{ 2r_D\frac{Q^2M^2}{(Q^2+M^2)^{(x)}2} + \frac{A_{3\text{IP}}}{b_{3\text{IP}}}\frac{M^2}{Q^2+M^2} \right\} =$$

$$\frac{R_{inc} + R_{coh}G_A(2m_N^2x^2[1+\exp(y)]^2)}{R_{sh}(x)} \cdot \left\{ 2r_D\frac{exp(y)}{[1+\exp(y)]^2} + \frac{A_{3\text{IP}}}{b_{3\text{IP}}}\frac{\exp(y)}{1+\exp(y)} \right\} \quad (32)$$

At small $x \ll 1/R_A m_N$, the major effect is an enhancement of large rapidity gaps $\Delta\eta \sim \log(\frac{1}{x})$, i.e., of excitation of $M^2 \sim Q^2$, by the factor

$$\frac{F_A(y)}{F_N(y)} = \frac{R_{coh} + R_{inc}}{R_{sh}(x)} > 1 \quad (33)$$

For the Xe nucleus, used as a target in the Fermilab E665 experiment [13], we predict an enhancement as strong as $F_A(y)/F_N(y) \sim 2.5$, see Fig. 3. This enhancement comes entirely from the coherent DD, and is partly due to $R_{sh}(x) < 1$.

Because of the form-factor suppression of the coherent DD cross section, the enhancement decreases with the increase of x, since the minimal longitudinal momentum transfer (7) increases with x. Similar suppression takes place with the increase of the mass of the excited state, i.e., with the increase of y and the decrease of the rapidity gap $\Delta\eta$. Because the form-factor suppression is absent in the incoherent DD, at large values of y the rapidity gap distribution for the nuclear target will be dominated by the incoherent production and

will flatten at the value which is smaller than for the free nucleon by the factor $R_{inc}/R_{sh} < 1$:

$$\frac{d\sigma_D^{(A)}}{\sigma_A d\Delta\eta} = \frac{A_{3\mathbb{P}}}{b_{3\mathbb{P}}} \frac{R_{inc}}{R_{sh}(x)} . \tag{34}$$

In Fig. 3 we show the y-distribition for the μXe interactions for few values of x assuming $W^2 = 400\text{GeV}^2$. As a matter of fact, in the kinematical range of the E665 experiment, the form factor effects are quite significant and there is no room for the flat y-distribution (34) at large y. Consider, for instance, $W^2 = 400\text{GeV}^2$ and $Q^2 = 2\text{GeV}^2$, i.e., $x = 5 \cdot 10^{-3}$. For the Xe nucleus the charge radius $R_{ch} \approx 4.5\text{f}$ [24]. The suppression by the form factor becomes significant at $R_{ch} m_N x[1 + \exp(y)] \gtrsim 1$, i.e., at $M^2/Q^2 \gtrsim 9$ and $y \gtrsim 2$. This corresponds to the onset of the triple-pomeron region, but still larger values of y are needed to reach the dominance of the incoherent DD. However, at this value of x, the requirement of $\Delta\eta \gtrsim 3$ imposes the upper bound $y \lesssim 2.3$, see Eq. (12). Even at $x = 2 \cdot 10^{-3}$ the triple-pomeron plateau is still elusive. At still larger values of x, the form factor effects become important at smaller values of y. This leads to the non-scaling $F_A(y)$ for DD on nuclei, whereas in DD on the free nucleons $F_N(y)$ is predicted to not depend on x and Q^2. The nuclear enhancement of large rapidity gaps, of small y and small M^2, and nuclear suppression of smaller rapidity gaps, i.e., of large y, is a very specific prediction, which can easily be tested experimentally. As we stated above, because of finite energy only the former prediction can be tested in the energy region of the E665 experiment.

In Fig. 5 we present the fraction of diffraction dissociation

$$W_D(y^*) = \int^{y^*} dy F(y) \tag{35}$$

as a function of y^* at different values of x. In Fig. 6 we present our prediction for the total, coherent plus incoherent, rate of DD $w_D(x)$ as a function of x for the free-nucleon and Xe targets for the rapidity-gap cutoff $\Delta\eta \gtrsim 3$. Bear in mind that the upper bound on y^* changes with x. At small x, the fraction of DD and/or LRG events on a nucleus is significantly higher than for the free-nucleon target. At large x, fractions of DD in μXe and μN interactions converge. Because of the nuclear form factor effects, nuclear coherent DD vanishes at $x \gtrsim 0.1$, and here a probability of LRG events in the μXe case even will be smaller than in the μN case. For the nucleon target, apart from the x-dependence of

the width of the allowed region of y, the probability $W_D(y)$ only weakly depends on x and Q^2. For the nuclear target there is a slight violation of this scaling, related to the scaling violation by the nuclear form factor effects in $F_A(y)$.

Notice, that whereas the diffraction dissociation is necessarily accompanied by the formation of a large (pseudo)rapidity gap, the reverse is not necessarily true. Specifically, even in the non-diffractive, valence dominance region of DIS, the virtual photon can fragment into the hadronic system of very small mass $M^2 \ll Q^2$, which will be separated from the recoil nucleon by large (pseudo)rapidity gap $\Delta\eta = \log(W^2/M^2)$. Such a non-diffractive rapidity-gap events must have an origin in the multiplicity fluctuations [30] and/or the secondary reggeon exchange across the rapidity gap, and must have the differential rapidity gap distribution which decreases exponentially at large rapidity gap,

$$F_{ND}(y) \propto \exp(-\Delta\eta) \propto x \cdot \exp(y) \tag{36}$$

compared to the flat or even rising probability of large diffractive gaps $\Delta\eta \sim \log(\frac{1}{x})$, i.e., of $y \sim 0$. Indeed, we predict quite a striking rise of $F_A(y)$ towards $y \to 0$, see Fig. 3.

The preliminary results from the E665 experiment on μXe interactions [31] confirm the above predictions. The E665 defines the large-rapidity gap (LRG) events subject to the psedurapidity gap $\Delta\eta > 2$, and separates the total statistics into the shadowing $x \leq 0.02$ and non-shadowing $x \geq 0.02$ samples. Their free-nucleon sample comes from the μD interactions. The E665 lower bound for the fraction of DD is 0.12 ± 0.02 for the μD and 0.18 ± 0.03 for the μXe interactions in the shadowing region [31], which is consistent with our prediction of nuclear enhancement of DD.

5.2 DD and grey tracks

The nuclear multiproduction events are conveniently classified according to the multiplicity n_g of the so-called grey tracks, which are predominantly recoil protons with the momentum $(150 - 200) \lesssim p \lesssim 600\,\mathrm{MeV/c}$. Here the lower cutoff is usually so chosen as to exclude the nucleus evaporation products. The multiplicity of grey tracks n_g measures the mutiplicity of inelastic intranuclear interactions. Still another similar observable is the total charge Q_T of secondary hadrons. In interaction with the free-nucleon (deuteron) target $\langle Q_T \rangle = \frac{1}{2}$.

The $Q_T \geq 2$ in multiproduction on a nucleus is an unambiguous signature of intranuclear cascading.

Evidently, the coherent DD on a nucleus entirely falls into the $n_g = 0$ and/or $Q_T = 0$ cathegory. In the incoherent diffraction dissociation the longitudinal recoil κ is smaller than the lower cutoff for the grey tracks. The transverse recoil momentum

$$p_\perp \sim b_D^{-1/2}, b_{3\mathbf{IP}}^{-1/2} \sim 300 - 400\,\mathrm{MeV/c} \tag{37}$$

can be sufficiently large to partly contribute to the $n_g \geq 1$, $Q_T \geq 1$ cathegory. However, as we have seen above, the incoherent DD only makes a small fraction of DD on a nucleus, see Fig. 1. Consequently, the above predicted nuclear emhancement of DD leads to a prediction of the decrease of $\langle n_g \rangle$, $\langle Q_T \rangle$ with the decrease of x from the nonshadowing (NS) region of $x \gtrsim 1/R_A m_N$ to the shadowing (Sh) region of $x \ll 1/R_A m_N$. To a crude approximation, we may assume that the multiplicity of grey tracks and/or the total observed charge in the non-diffractive interactions do not depend on x , which leads to the estimate

$$\langle n_g \rangle_{Sh} \approx [1 - W_D(x)]\langle n_g \rangle_{NS}\,,$$

$$\langle Q_T \rangle_{Sh} \approx [1 - W_D(x)]\langle Q_T \rangle_{NS}\,. \tag{38}$$

Our estimate for the suppression factor $1 - W_D(x)$ is shown in Fig. 7. Notice, that in the hadron-nucleus interaction $\langle n_g \rangle$ rises with $\bar\nu$, which by the simple-minded extrapolation would have suggested the rise of $\langle n_g \rangle$ from the non-shadowing region of large x to the shadowing region of small x [1,2,8-10]. As a matter of fact, this study was primarily motivated by the preliminary evidence for such a x-dependence of $\langle n_g \rangle$ in the E665 data on μXe interaction [31]. The E665 data show an $\approx 30\%$ reduction of $\langle n_g \rangle$ and Q_T from $x \gtrsim 0.1$ to $x \sim 0.001$, which is in very good agreement with the estimate (38) shown in Fig. 7. One would expect [8-10] stronger intranuclear cascading and certain enhancement of $\langle n_g \rangle$ going from the non-shadowing to the shadowing region of x, so that Eq. (38) gives rather the lower bound for the suppression factor.

The fraction of diffractive production $W_D(x)$ is a scaling function of x, and we advocate binning the experimental data *vs.* x rather than *vs.* Q^2 and/or W^2. Small values of x are only accessible at high energy ν, and we predict a decrease of $\langle n_B \rangle$, $\langle n_g \rangle$ and $\langle Q_T \rangle$ with

increasing energy.

5.3 DD and the (pseudo)rapidity spectrum of secondary particles

In the non-diffractive (ND) inelastic interaction, secondary particles populate the whole rapidity span. By the formation-time considerations, the forward particle production in the ND events must be target independent [1,8-10]. In the DD events the secondary paricles populate only the photon hemisphere, with the vanishing activity in the target region for the coherent DD, and with some signal from the recoil protons in the incoherent DD. Now we comment in more detail on the signal of nuclear enhancement of DD in the (pseudo)rapidity spectra. The small effects of DD on the very forward hadroproduction on nuclei were discussed earlier in [18].

A convenient quantity is the normalized (pseudo)rapidity η distribution of secondary particles in the photon-proton center-of-mass system

$$R(\eta) = \left(\frac{dn^{(A)}}{d\eta} \right) \left(\frac{dn^{(N)}}{d\eta} \right)^{-1},$$

and let $\eta > 0$ be the photon hemisphere. In the generic particle-nucleus interaction $R(\eta)$ rises towards the target fragmentation region because of the nuclear cascading [1,2,8-12]. In the hadron-nucleus interactions, because of simultaneous interactions of few constituent quarks of the hadron, in the central region $R(\eta) > 1$, and in the projectile fragmentation region $R(\eta) < 1$ (for the review see [1]). In the leptoproduction on nuclei in the photon fragmentation region $R(y^*) \approx 1$ with some evidence for nuclear depletion in the maximal-y^* bin [12,13].

Evidently, the nuclear enhancement of DD leads to a higher particle density in the forward hemisphere in the multiproduction on nuclei compared to the free nucleon (deuterium) target. However, this simple prediction is not easy to test. Firstly, the total rapidity span depends on W^2. Because the rapidity spectrum $dn^{(N)}/d\eta$ is a steep function of η, small mismatch in W^2 may lead to large spurious effects in $R(\eta)$. Secondly, the diffraction dissociation products are smeared over broad rapidity range $\propto \log(M^2)$ and such a smearing is even broader in the pseudorapidity variable. The effect of the smearing tends to diminish

the departure of $R(\eta)$ from unity [1]. Very crude estimate for forward production is

$$R_{A/D}(\eta) - 1 \approx W_D^{(A)}(x) - W_D^{(N)}(x) \tag{39}$$

Notice, that because of the nuclear suppression of DD in the nonshadowing region of large x, in the non-shadowing region we expect $R_{A/D}(\eta) < 1$ for the very forward particles.

One may also compare the nuclear interactions in the non-shadowing and shadowing region, where one would expect

$$R_{sh/NS}(\eta) - 1 \approx W_D^{(A)}(x) \tag{40}$$

Here one compares the spectra at different values of Q^2, and the effect can be masked by weak Q^2 dependence of the rapidity spectra. The effect can somewhat be enhanced, if one compares the forward pseudorapidity spectra in the $n_g = 0$ events in the shadowing and non-shadowing regions on the same nucleus. The sample of the $n_g = 0$ events will evidently be enriched by DD. If $P_{NS}(0)$ is the probability of having $n_g = 0$ for the nonshadowing region of large x, then in the shadowing region we expect

$$P_{sh}(0) \approx P_{NS}(0) + w_D(x) \,. \tag{41}$$

Whith the above reservations about possible Q^2 dependence of the spectra, we expect

$$R_{sh/NS}(n_g = 0, \eta) \approx 1 + \frac{w_D(x)}{P_{NS}(0)} \tag{42}$$

Similar effect is expected, if one compares the forward production in the $n_g = 0$ and the $n_g \geq 1$ samples.

In all the above cases, this enhancement can, perhaps, best be seen by a comparison of average multiplicities in the forward hemisphere $\langle n_F \rangle$. For instance, in view of Eq. (41) we predict that $\langle n_F \rangle$ in the $n_g = 0$ sample of μXe interactions must be larger than for the free-nucleon target. We expect a similar enhancement of the forward multiplicity in the shadowing region of small x compared to the non-shadowing region of large x, although such a comparison is somewhat indirect because of different values of Q^2 in the two regions and possible slight dependence of mean multiplicty on Q^2. Similar enrichement by diffraction dissociation must hold also for the forward production in the $Q_T = 0$ sample of nuclear interactions.

Consider now the multiplicities in the target hemisphere. Here the intranuclear cascading leads to $R_{(\eta)} > 1$ in the global rapidity spectra. Because DD does not contribute to the target hemisphere, the rising fraction $w_D(x)$ of DD must be associated with the decrease of the mean multiplicity $\langle n_B \rangle$ of secondary partilces produced in the nucleus hemisphere, cf. Eq. (38):

$$\frac{\langle n_B \rangle_{Sh}}{\langle n_B \rangle_{NS}} \approx 1 - w_D(x) \tag{43}$$

This estimate is in very good agreement with the preliminary data from the E665 experiment [31], which found $\approx 30\%$ depletion of $\langle n_B \rangle$ from $x \gtrsim 0.1$ to $x \sim 0.001$, and the observed depletion is the same as for $\langle n_g \rangle$ and Q_T.

6 What makes the hadroproduction and leptoproduction on nuclei different?

Diffraction dissociation and LRG interactions exist in the hadronic interactions too, but their effect on the multiproduction on nuclei is marginal. The principal distinction between the leptoproduction and hadroproduction is strong absorption via elastic rescatterings in the latter case. Furthermore, the diffraction dissociation of hadrons makes only a small fraction of diffractive scattering, which is dominated by elastic scattering.

Let us make this argument more explicit. The formalism of Section 2 is fully applicable to diffractive scattering of hadrons. The total cross section of hN scattering can be written as $\sigma_{tot}(hN) = \langle \hat{\sigma} \rangle_h$, where the subscript h denotes the matrix element (5) over the wave function of the hadron h and $\hat{\sigma}$ stands for the generic cross section operator, $\hat{\sigma} = \sigma(r)$ in the example considered in Section 2. The diffrential cross section of the forward elastic scattering can then be written as (we suppress the subscript h)

$$\left. \frac{d\sigma_{el}}{dt} \right|_{t=0} = \frac{1}{16\pi} \langle \hat{\sigma} \rangle^2 \,. \tag{44}$$

If all eigenvalues of the diffraction matrix were identical, then the diffractively scattered wave would have differed from the incoming hadronic wave only by the overall phase/attenuation

factor, and there would not have been any diffraction dissociation (for the review see [1]). The counterpart of Eq. (6) for hadrons is [1,22]

$$\frac{d\sigma_{el}}{dt}\bigg|_{t=0} = \frac{1}{16\pi} \left(\langle \hat{\sigma}^2 \rangle - \langle \hat{\sigma} \rangle^2 \right) \tag{45}$$

and DD measures the dispersion of eigenvalues of the diffraction scattering operator. For the proton-proton scattering, the detailed analysis of the DD data gives [32]

$$\frac{\langle \hat{\sigma}^2 \rangle - \langle \hat{\sigma} \rangle^2}{\langle \hat{\sigma} \rangle^2} \sim 0.3 \,. \tag{46}$$

At high energy, elastic scattering makes $\sim 20\%$ of the total pp cross section, and single-arm DD makes only $\sim 6\%$ of the total cross section. In DIS at small x we found much larger fraction of DD, $\sim 15\%$ in the μN scattering and $\sim (35 - 40)\%$ in the μXe scattering at $x = 0.002$. Is this reasonable?

In the case of photons the term $\langle \hat{\sigma} \rangle^2$ is negligibly small, as it contains the extra power of the fine structure constant $\alpha_{em} = 1/137$. Then, the comparison of Eq. (6) with Eqs. (45,46) shows that the strength of DD in the photoabsorption corresponds to the combined strength of elastic scattering and the beam DD in the hadronic scattering. This explains why we find such a strong DD for photons. Furthermore, consider the total photoabsorption cross section Eq. (15) and the coherent DD cross section Eq. (19) in the limit of very heavy nucleus, when the absorption becomes strong. In this limit one will find that the coherent DD on a nucleus makes $\frac{1}{2}$ of the total cross section,

$$\sigma_{coh} \sim \frac{1}{2}\sigma_A \,, \tag{47}$$

which precisely corresponds to the black disc limit when the elastic cross section equals half of the total cross section.

The above rise of the coherent DD cross section to $\sim \frac{1}{2}$ of the total DIS cross section must be contrasted with very small cross of DD of hadrons on nuclei, which simply vanishes in the black-disc limit. The principal point is the following one. The above estimate (46) shows that dispersion of the cross section for Fock states of the hadron is not large. Then, in the hadron-nucleus scattering, the overall nuclear attenuation will be dominated by the average value of the cross section, i.e., by the free-nucleon cross section. In other words, it

will be dominated by elastic rescatterings of the projectile hadron, and the Glauber formula for the nuclear total cross section [3]

$$\sigma_{hA} = 2 \int d^2\vec{b} \{1 - \exp[-\frac{1}{2}\sigma_{hN}T(b)]\} \qquad (48)$$

gives an excellent description of the nuclear shadowing, which in the nA scattering on heavy nuclei reduces the total cross section by more than the factor 2 (for analysis of nuclear shadowing in hadron-nucleus scattering see [7]). The diffraction dissociation effects are present in this case too and contribute to the nuclear shadowing. The corresponding correction to the nuclear shadowing, usually referred to as Gribov's inelastic shadwoing, can be evaluated as [1,4,7,24]

$$\Delta\sigma_{sh} = 4\pi \left.\frac{d\sigma_D^{(N)}}{dt}\right|_{t=0} \int d^2\vec{b} \, T(b)^2 \exp[-\frac{1}{2}\sigma_{hN}T(b)] + ... \qquad (49)$$

and only makes 2-4% correction to the impulse approximation cross section and 5-7% correction to the total nA cross section [5,22]. The cross section of the diffraction dissociation on nuclei will be even smaller because of the stronger nuclear attenuation factor:

$$\sigma_D = 4\pi \left.\frac{d\sigma_D^{(N)}}{dt}\right|_{t=0} \int d^2\vec{b} \, T(b)^2 \exp[-\sigma_{hN}T(b)] + ... \qquad (50)$$

For this reason, LRG events only contribute a negligible fraction to the hadronic multiproduction on nuclei, and their effect is only noticable for very fast particle production [18]. Significant nuclear enhancement of the mean multiplicity in the hadroproduction on nuclei comes about equally from reinteractions of the projectile and from the cascading effects [33,34].

7 Conclusions

Our principal finding is a strong nuclear enhancement of the diffraction dissociation of (virtual) photons in deep inelastic leptoproduction off nuclei. This enhancement of diffraction dissociation amounts to enhancement of the large-rapidity gap and/or small multiplicity events in the nuclear shadowing region, in which photons are expected to have stronger hadronic properties. The nuclear enhancement of the coherent diffraction dissociation also

amounts to suppression of the target-nucleus fragmentation. Henceforth, a somewhat para-doxial conclusion that the more hadronlike photons produce less particles on nuclei. The principal difference from the hadron-nucleus interactions is that in the hA case the nuclear shadowing is dominated by the elastic rescatterings of the projectile hadron, and that the diffraction dissociation effects, alias the inelastic shadowing, only make a very small correction to the nuclear cross section. Diffraction dissociation of hadrons vanishes for the black nuclei. In contrast to that, in deep inelastic scattering on the completely absorbibg, black, nuclei the coherent diffraction dissociation must make $\sim \frac{1}{2}$ of the total nuclear DIS cross section.

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Figure captions:

Fig.1 - The A-dependence of the contribution from interaction of the $q\bar{q}$ Fock state of the photon to the DD related quantities: nuclear shadowing R_{sh} is shown by the dotted curve; the nuclear enhancement of the coherent DD cross section R_{coh} and suppression of the incoherent DD cross section R_{inc} are shown by the dashed and dot-dashed curve; the solid curve is for the ratio of the total DD cross section to the shadowing cross section. The definitions are given in the text Eqs. (15,19,20,23).

Fig.2 - Nuclear shadowing in μXe scattering as a function of x. The dashed line shows the contribution to shadowing from the $q\bar{q}$ Fock state of the photon, the solid line includes also the triple-pomeron component from shadowing of higher Fock states.

Fig.3 - The diffraction dissociation mass spectrum Eqs. (29,32) as a function of the scaling variable $y = \log(M^2/Q^2)$ The solid curve is for μN interaction at $x = 0.002$ the dashed, dotted and dash-dotted curves are for μXe interaction at $W^2 = 400 \mathrm{GeV}^2$ and different values of x.

Fig.4 - The diffraction dissociation mass spectrum in μN interaction at $W^2 = 400 \mathrm{GeV}^2$ and different values of x is shown as a function of M^2.

Fig.5 - The fraction of diffraction dissociation Eq. (35) integrated over excited masses $M^2 \leq Q^2 \exp(y^*)$ at $W^2 = 400 \mathrm{GeV}^2$ and different values of x. The end points of curves correspond to the rapidity gap $\Delta\eta = 3$.

Fig.6 - The x-dependence of the fraction of diffraction dissociation integrated over rapidity gaps $\Delta\eta \geq 3$ at $W^2 = 400 \mathrm{GeV}^2$.

Fig.7 - The estimate of suppression of grey particle multiplicity because of nuclear enhancement of DD.

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